

Electric dipole moments of charged leptons at one loop in presence of massive neutrinos

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Violation of CP invariance is a quite relevant phenomenon that is found in the Standard Model, though in small amount. This has been an incentive to look for high-energy descriptions in which CP violation is increased, thus enhancing effects that are suppressed in the Standard Model, such as the electric dipole moments of elementary particles. In the present investigation, we point out that charged currents in which axial couplings are different from vector couplings are able to produce one-loop contributions to electric dipole moments of charged leptons if neutrinos are massive and if these currents violate CP . We develop our discussion around charged currents involving heavy neutrinos and a W' gauge boson coupling to Standard Model charged leptons. Using the most stringent bound on the electron electric dipole moment, provided by the ACME Collaboration, we determine that the difference between axial and vector currents lies within $\sim 10^{-10}$ and $\sim 10^{-11}$ for heavy-neutrino masses between $\sim 10^2$ GeV and $\sim 10^3$ GeV. This possibility is analyzed altogether with the contributions to the anomalous magnetic moments of charged leptons, which establish strict requirements, including a quasi-degenerate spectrum of heavy-neutrino masses. In this context, we estimate, for the τ lepton, an anomalous magnetic moment between $\sim 10^{-8}$ and $\sim 10^{-10}$ for neutrino masses ranging from 0.5 TeV to 7 TeV.

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I. INTRODUCTION

The discovery of the Higgs-like particle [1, 2] with mass around 125 GeV, announced by the CMS [3] and ATLAS [4] Collaborations at the Large Hadron Collider, has been a remarkable achievement, which, however, is not a statement that the Standard Model is the last part of the story. The premise that there is a more fundamental physical description, beyond the Standard Model, was fed during years by theoretical issues, but hints of its nature have been finally provided by experimental observations that include neutrino oscillations [5, 6], dark matter [7–9], and perhaps even a new particle with mass ~ 750 GeV [10, 11]. The phenomenon of neutrino oscillations, first observed at Super-Kamiokande, then at the SNO, and recently confirmed by the determination of the last mixing angle by the Daya Bay [12] and RENO [13] Collaborations, has been interpreted as an effect of neutrino mixing and neutrino mass [14, 15]. Among other things, this event set the quite relevant question of whether the neutrinos correspond to Dirac or Majorana fermions. Clues on the answer might come from experimental searches of the elusive neutrinoless double beta decay. It has been pointed out that the electromagnetic properties of massive neutrinos are very different depending on whether these fermions are of Dirac or Majorana type [16–19], but they are elusive and difficult to analyze. An important aspect of neutrino mixing is that the measurement of a nonzero value of the θ_{13} mixing angle rendered it a source of CP violation. The Pontecorvo-Maki-Nakagawa-Sakata mixing matrix [20, 21] is able to introduce violations

of such invariance by means of one complex phase, if neutrinos are Dirac fermions, or even three phases, in case that neutrinos are Majorana fermions [15, 22].

Certainly, searches for deviations from Standard Model predictions deserve much attention. The exploration of processes that are quite suppressed, or even forbidden, in the Standard Model may eventually find hints about some theory describing nature beyond this low-energy description. According to Sakharov criteria, the nonconservation of CP invariance is a necessary requirement for the baryon asymmetry to occur [23]. The violation of CP symmetry is indeed an effect that is included in the Standard Model, though in small amount, by the complex phase of the Cabibbo-Kobayashi-Maskawa matrix [24, 25]. Studies aimed at other sources of CP -violation, from physics beyond the Standard Model, constitute an active topic nowadays. In particular, diverse investigations as, for instance, those performed in Refs. [26–37], have explored the generation of electric dipole moments of elementary particles through CP -violation characterizing new-physics formulations. Among all electric dipole moments of elementary particles, that of the electron is, doubtless, the one which has been most stringently bounded [38, 39].

Violation of CP invariance in neutrino mixing might induce electric dipole moments of charged leptons. We explore, in the present paper, the impact of lepton general charged currents on the electromagnetic form factors of Standard Model charged leptons at one loop. The charged currents that we consider involve Standard

Model charged leptons, l_α , a heavy charged gauge boson, W' , and a set of heavy Dirac neutrinos, N_j . We find that electric dipole moments of charged leptons arise at the one-loop level if three conditions hold: 1) the charged currents violate CP ; 2) the axial and vector terms in such currents differ of each other; and 3) the neutrinos are massive. This is in contrast with the Standard Model contributions from the Cabibbo-Kobayashi-Maskawa phase to the electric dipole moment of the electron, which vanish even at the three-loop order [40] and produce, at four loops, the tiny value $\sim \mathcal{O}(10^{-44}) e \cdot \text{cm}$ [41].

Anomalous magnetic moments are connected with the aforementioned contributions and they establish stringent constraints on them. The masses $m_{W'}$ and m_j , of the W' boson and the heavy neutrinos N_j , are assumed to grow with some high-energy scale Λ . As we show below, the contributions from any neutrino N_j to any anomalous magnetic moment decouple as $\Lambda \rightarrow \infty$. Furthermore, the control on the size of these contributions, and consequently on whether they actually make sense on experimental grounds, is set by how similar these masses are: if $m_{W'}$ and m_j are very different of each other, the contributions are anomalously large. This requirement also implies that the heavy-neutrino mass spectrum is quasi-degenerate. In this context, we analyze which regions are under control and how the corresponding values affect the contributions to the electric dipole moments. Being the difference among vector and axial charged currents a necessary condition to produce one-loop electric dipole moments, we use the upper limit, of order $10^{-29} e \cdot \text{cm}$, on the electric dipole moment of the electron [38, 39] to estimate an upper bound $\sim 10^{-11} - 10^{-10}$ on such difference for the case of heavy neutrinos, with masses in the TeV range. We also find a contribution within $10^{-10} - 10^{-8}$ to the anomalous magnetic moment of the tau lepton.

The paper has been organized in the following manner: in Section II, we define our framework and sketch the calculation of the one-loop contributions to the electromagnetic vertex; in Section III, we explore the contributions to magnetic moments, we analyze the restrictions that are imposed by these quantities, and we calculate the contribution to the tau anomalous magnetic moment; Section IV is dedicated to the one-loop electric dipole moments contributions, and the upper bound on the difference among vector and axial currents is derived for the case of the electron; finally, Section V is used to present our conclusions.

II. ELECTROMAGNETIC MOMENTS FROM GENERAL CHARGED CURRENTS

We start by considering the general set of charged currents (CC)

$$\mathcal{L}_{\text{CC}} = \frac{1}{2\sqrt{2}} \sum_j \sum_\alpha \left[W_\rho'^+ \bar{N}_j \gamma^\rho (v_{j\alpha} - a_{j\alpha} \gamma^5) l_\alpha + \text{H.c.} \right], \quad (1)$$

where $\alpha = e, \mu, \tau$ is a flavor index, so that l_α represents charged leptons, and $j = 1, 2, 3$ runs over heavy Dirac neutrinos N_j . We are assuming that the axial and vector couplings are different of each other, that is, $v_{j\alpha} \neq a_{j\alpha}$. The set of coefficients $v_{j\alpha}$ and $a_{j\alpha}$ implicitly bear all the information about heavy-neutrino mixing, which we assume to violate CP invariance. While we are restricting our study to heavy neutrinos, note that, in a more general context, the sum over j could involve both heavy and light neutrinos. In such case, the set of coefficients $v_{j\alpha}$ and $a_{j\alpha}$, in Eq. (1), could be viewed as entries of nonsquare complex matrices, which thus would not be restricted to be unitary, as it is the case of the neutrino mass model analyzed in Ref. [42]. The greek index ρ , in the W' charged boson field, labels spacetime coordinates, with a sum over any pair of repeated indices. A more general set of charged currents could include even more charged gauge bosons, as those considered in Ref. [43] to calculate Majorana neutrino magnetic moments in *left-right models* [44].

The one-loop contributions from the charged currents given in Eq. (1) to the electromagnetic vertex $\gamma l_\alpha l_\beta$, with l_α and l_β being either equal or different, emerge from the Feynman diagrams shown in Fig. 1. We perform this

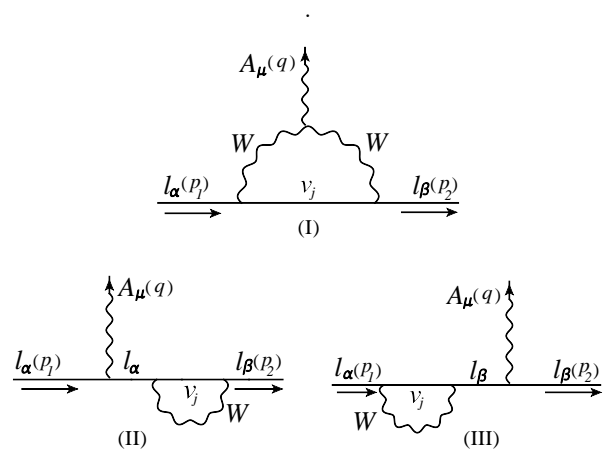


FIG. 1: One-loop diagrams contributing to the flavor-changing electromagnetic vertex.

calculation in the unitary gauge. Taking the external particles on shell, we find that the resulting vertex function has the well-known structure of the electromagnetic vertex parametrization, which is, for instance, provided

in Refs. [27, 45] and which reads¹

$$\Gamma_{\mu}^{\alpha\beta} = ie \left[\gamma_{\mu} (f_{\alpha\beta}^V - f_{\alpha\beta}^A \gamma_5) - \sigma_{\mu\nu} q^{\nu} \left(i \frac{\mu_{\alpha\beta}}{m_{\alpha} + m_{\beta}} - \frac{d_{\alpha\beta}}{e} \gamma_5 \right) \right]. \quad (2)$$

Here, e denotes the unit electric charge, and m_{α} and m_{β} are the masses of external charged leptons l_{α} and l_{β} . The parameters $f_{\alpha\beta}^V$, $f_{\alpha\beta}^A$, $\mu_{\alpha\beta}$, and $d_{\alpha\beta}$ are, respectively, the one-loop contributions to electric charge, axial current, anomalous magnetic, and electric dipole moments. Even though this calculation was carried out on shell and in a specific gauge, all the factors in Eq. (2) are complicated functions of the masses of all the fields involved in the contributing diagrams (see Fig. 1). These are the W' boson mass $m_{W'}$, the masses m_{α} and m_{β} of charged leptons, and the N_j neutrino masses m_j .

We have verified that for $\alpha \neq \beta$ the contributions to $f_{\alpha\beta}^V$ and $f_{\alpha\beta}^A$ vanish exactly. If $\alpha = \beta$ the corresponding nonzero contributions contain ultraviolet divergences, but they are expected to be absorbed by renormalization. The factors $\mu_{\alpha\beta}$ and $d_{\alpha\beta}$ corresponding to $\alpha \neq \beta$ are respectively called *transition magnetic moments* and *transition electric moments*. On the other hand, the $\mu_{\alpha} \equiv \mu_{\alpha\alpha}$ are the *anomalous magnetic moments* and the $d_{\alpha} \equiv d_{\alpha\alpha}$ are the *electric dipole moments*. To calculate these loop contributions we have utilized the Passarino-Veltman method [48], which means that all these electromagnetic moments are expressed in terms of two-point and three-point scalar functions, B_0 and C_0 . While the C_0 functions are ultraviolet finite, the B_0 functions include such type of divergences. In the dimensional regularization approach [49] any B_0 function can be expressed as [50] $B_0 = \Delta_{\text{div}} + f_{\text{fin}}$, with all the ultraviolet divergences and the logarithmic cutoff dependence contained in Δ_{div} , which is shared by all the two-point functions. Taking advantage of this generic form of the B_0 functions, we have checked that all ultraviolet divergences in both transition and diagonal magnetic and electric moments are exactly eliminated, yielding finite results for these quantities. This is consistent because the magnetic and electric dipole Lorentz structures are exclusively generated at the loop level.

Our calculation reveals that the general structure of

the magnetic and electric moments is

$$\mu_{\alpha} = \sum_j (|a_{j\alpha}|^2 M_{j\alpha}^A + |v_{j\alpha}|^2 M_{j\alpha}^V), \quad (3)$$

$$d_{\alpha} = \sum_j i (a_{j\alpha}^* v_{j\alpha} - v_{j\alpha}^* a_{j\alpha}) D_{j\alpha} \cdot e, \quad (4)$$

$$\mu_{\alpha\beta} = \sum_j [a_{j\alpha}^* a_{j\beta} M_{j,\alpha\beta}^A + v_{j\alpha}^* v_{j\beta} M_{j,\alpha\beta}^V], \quad (5)$$

$$d_{\alpha\beta} = \sum_j i [v_{j\alpha}^* a_{j\beta} D_{j,\alpha\beta} + a_{j\alpha}^* v_{j\beta} \overline{D}_{j,\alpha\beta}] \cdot e, \quad (6)$$

each of them with a sum over the index j , because of neutrinos circulating in the loops. All the dependence of these electromagnetic moments on Passarino-Veltman functions and, in general, on the masses of particles (see Fig. 1), lies within the real-valued factors $M_{j\alpha}^A$, $M_{j\alpha}^V$, $D_{j\alpha}$, $M_{j,\alpha\beta}^A$, $M_{j,\alpha\beta}^V$, $D_{j,\alpha\beta}$, and $\overline{D}_{j,\alpha\beta}$. The explicit expressions of all these factors, in terms of Passarino-Veltman scalar functions, are provided in Appendix A. Eq. (3) shows that anomalous magnetic moments μ_{α} are real quantities, which also occurs with electric dipole moments d_{α} , Eq. (4), whose terms within the neutrino sum are proportional to $\Im(v_{j\alpha}^* a_{j\alpha})$. The transition moments $\mu_{\alpha\beta}$ and $d_{\alpha\beta}$ are, in general, complex numbers. Let us point out that the charged currents provided in Eq. (1) violate CP invariance if at least $v_{j\alpha}$ or $a_{j\alpha}$ is a complex number, but they preserve this symmetry if both of them are real quantities [43]. As it can be appreciated from Eq. (4), contributions to the electric dipole moment of a charged lepton l_{α} are nonzero only if CP is violated by the charged currents of Eq. (1) and $v_{j\alpha} \neq \pm a_{j\alpha}$. Eqs. (3), (5), and (6) clearly show that conservation of the CP symmetry and/or fulfillment of the condition $v_{j\alpha} = \pm a_{j\alpha}$ does not forbid the existence of contributions to all other electromagnetic moments. In particular, the status of CP symmetry, in this context, is completely irrelevant to the anomalous magnetic moment μ_{α} , as it can be observed in Eq. (3).

III. CONSTRAINTS FROM MAGNETIC MOMENTS

The diagonal magnetic moments that we just showed, in Eq. (3), are sums of contributions from massive heavy neutrinos N_j . For each of such contributions we use the notation $\mu_{j\alpha}$, so that $\mu_{\alpha} = \sum_j \mu_{j\alpha}$. In Ref. [51], the Higgs decay into two quarks, in the context of the Standard Model, was recently revisited. An interesting element of this study was a “heavy mass limit”, in which the masses of internal quarks, $m_{q_{\text{int}}}$, and the W boson mass, m_W , were taken to be the same for a very large electroweak scale v , that is, $m_W = m_{q_{\text{int}}}$ for very large v . It was then observed and discussed that, as long as this condition is fulfilled, the decay amplitude $H \rightarrow q_i q_j$ goes

¹ As it is shown in Refs. [46, 47], violation of Lorentz invariance allows a richer structure of this parametrization.

to zero for $v \rightarrow \infty$, which was used as a consistency check.

In the present investigation, we notice that the same happens with any heavy-neutrino contribution $\mu_{j\alpha}$: if we take $m_{W'} = m_j$ and then the limit of $\mu_{j\alpha}$ as $m_j \rightarrow \infty$ is considered, it occurs that $\mu_{j\alpha} \rightarrow 0$. With this in mind, we assume that some high-energy scale, Λ , is associated with the generation of the masses m_j and $m_{W'}$, by spontaneous symmetry breaking. We assume that $m_{W'}$ and m_j grow with Λ and that $m_{W'} \rightarrow m_j$ as $\Lambda \rightarrow \infty$. Then we set the condition $m_{W'} = (1 - \epsilon_j) m_j$, where ϵ_j is given by

$$\epsilon_j = \frac{m_j - m_{W'}}{m_j}. \quad (7)$$

We wish to emphasize that each ϵ_j is independent of the lepton flavor that corresponds to the each anomalous magnetic moment. Evidently, the sign of ϵ_j depends on which of the heavy masses $m_{W'}$, m_j is the largest, with $\epsilon_j > 0$ if $m_j > m_{W'}$, but with $\epsilon_j < 0$ holding if $m_{W'} > m_j$. From Eq. (3), note that $|\mu_{j\alpha}| \leq |M_{j\alpha}^A + M_{j\alpha}^V|$. Now we assume that the parameter ϵ_j is be very small, in which case we have

$$|\mu_{j\alpha}| \lesssim \frac{1}{128\pi^2} \left| -\frac{1}{4} \frac{m_\alpha^2}{m_j^2} \left(\frac{m_\alpha^2}{m_j^2} - 21 \right) + \epsilon_j \left(-\frac{5}{2} + \frac{29}{3} \frac{m_\alpha^2}{m_j^2} - \frac{5}{6} \frac{m_\alpha^4}{m_j^4} \right) \right|. \quad (8)$$

This expression shows that any $\mu_{j\alpha}$ decouples as $m_j \rightarrow \infty$ ($\Lambda \rightarrow \infty$).

Now we analyze the contributions from the new physics, producing the charged currents of Eq. (1), to the anomalous magnetic moments of Standard Model charged leptons. Keeping in mind our discussion of the previous paragraph, it shall be clear in a moment that an appropriate estimation of the contributions to the anomalous magnetic moments has to take into account the size of the parameter ϵ_j . Besides taking a small ϵ_j , we introduce the further assumption that $|a_{j\alpha} - v_{j\alpha}| \ll m_\alpha/m_j$, which will make sense in the next section, when we discuss electric dipole moments. Under these circumstances, we have:

$$\mu_{j\alpha} \approx \frac{(|a_{j\alpha}|^2 + |v_{j\alpha}|^2)}{32(4\pi)^2} \left(\frac{31}{2} \frac{m_\alpha^2}{m_j^2} - 5\epsilon_j \right). \quad (9)$$

If ϵ_j is not small enough, the term $-5\epsilon_j$ in the right-hand side of Eq. (9), which is not suppressed by multiplicative powers of small charged-lepton masses m_α nor by inverse powers of large neutrino masses m_j , will be dominant. Indeed, for a large-enough ϵ_j the contributions to the anomalous magnetic moment of any charged lepton would be practically independent of the flavor of the lepton taken into consideration.

Among the anomalous magnetic moments of charged leptons, the most accurate investigations, in the context of the Standard Model, have been performed in the cases of the muon and the electron [52–54]. The current discrepancies among the values measured by experiments [55–57] and the predictions from the Standard Model are $\Delta\mu_\mu = \mu_\mu^{\text{exp}} - \mu_\mu^{\text{SM}} = 2.49 \times 10^{-9}$ and $\Delta\mu_e = \mu_e^{\text{exp}} - \mu_e^{\text{SM}} = -1.06 \times 10^{-12}$ [52, 53], which means that this is a good place to look for suppressed new physics. For the next step, we stress that, according to the analysis of the present paper, the mass of any heavy neutrino N_j cannot be far from the W' mass. Thus, in this model a quasi-degenerate heavy-neutrino mass spectrum is required. With this in mind, we consider some mass m_N and some ϵ_N such that $m_j \approx m_N$ and $\epsilon_j \approx \epsilon_N$ for any N_j , so that $\mu_\alpha = \sum_j \mu_{j\alpha} \approx 3\mu_{N\alpha}$. Using Eq. (9) and taking $|a_{N\alpha}|^2 + |v_{N\alpha}|^2 \sim 1$, we derive, from $\Delta\mu_e$ and $\Delta\mu_\mu$, a couple of regions where ϵ_N may lie, depending on the heavy mass m_N :

$$1.03 \left(\frac{m_e^2}{m_N^2} \right) < \epsilon_N < 1.03 \left(\frac{m_e^2}{m_N^2} \right) + 3.57 \times 10^{-10}, \quad (10)$$

$$1.03 \left(\frac{m_\mu^2}{m_N^2} \right) - 8.4 \times 10^{-7} < \epsilon_N < 1.03 \left(\frac{m_\mu^2}{m_N^2} \right). \quad (11)$$

For m_N between 0.5 TeV and 7 TeV, the factor $1.03 (m_e^2/m_N^2)$ ranges between $\sim 10^{-15}$ and $\sim 10^{-12}$, so that Eq. (10) sets an upper limit on ϵ_N that is practically determined by the constant term of order 10^{-10} . Because $1.07 (m_\mu^2/m_N^2)$ varies from 10^{-10} to 10^{-8} , for m_N between the same range of masses, something similar happens with the lower bound of Eq. (11), which turns out to be a negative number. Both allowed regions in the (m_N, ϵ_N) plane are shown in the upper density plot of Fig. 2, with the upper graph corresponding to Eq. (10) and the lower one to Eq. (11). How large are the values of μ_α for the different values of m_N and ϵ_N is represented in each case by a gradient, which goes from red (smaller values) to yellow (larger values). Recall that ϵ_N is independent of the lepton flavor, so that the actual allowed region for this parameter is the intersection of Eqs. (10) and (11), which corresponds to values $\epsilon_N \sim 10^{-13} - 10^{-10}$. With the allowed region for ϵ_N in our hands, we note that $m_\tau^2/m_N^2 \sim 10^{-8}$ for $m_N = 7$ TeV, which means that in the case of the tau lepton, the contribution of Eq. (9) is completely dominated by the first term within brackets. Thus we estimate, in this model, the value of the anomalous magnetic moment of the τ lepton, which we find to be $\mu_\tau \sim 10^{-8}$ if $m_N = 0.5$ TeV and $\mu_\tau \sim 10^{-10}$ for $m_N = 7$ TeV.

The Standard Model contribution to the anomalous magnetic moment of the τ lepton was accurately calculated in Ref. [58], where the value $\mu_\tau^{\text{SM}} = 117721(5) \times 10^{-8}$ was reported. This value is larger, by several orders of magnitude, than the prediction of the present work, which, however, is not so different from contributions

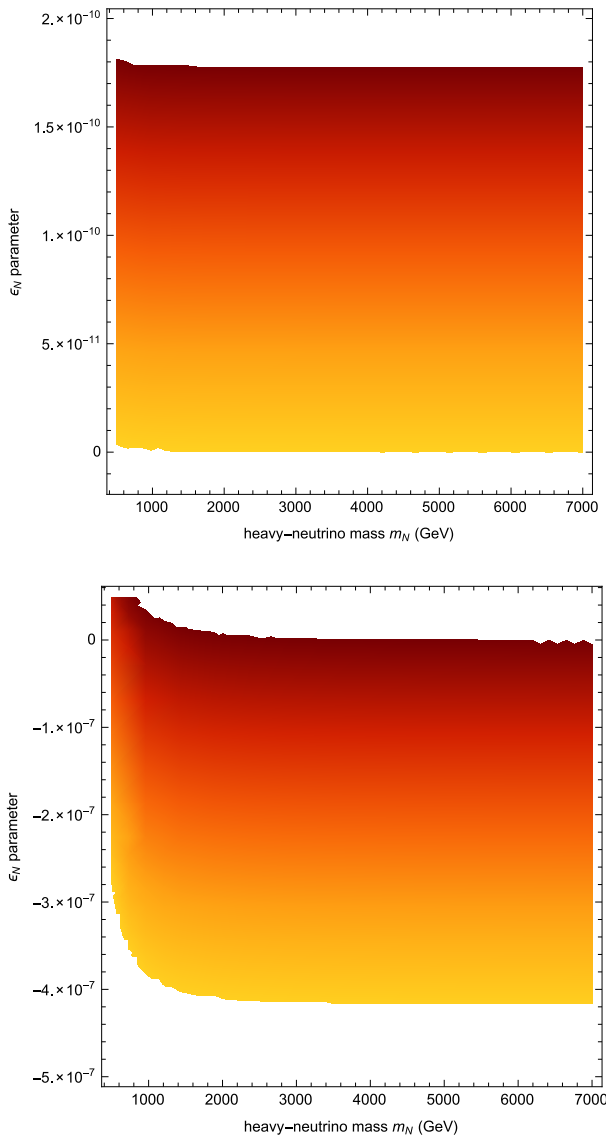


FIG. 2: Allowed regions of the (ϵ_N, m_N) plane for the $\Delta\mu_e$ difference (upper graph) and for the $\Delta\mu_\mu$ difference (lower graph).

produced by Standard Model extensions. In the case of the Minimal Supersymmetric Standard Model with a mirror fourth generation, the authors of Ref. [59] calculated contributions to μ_τ that range within $\sim 10^{-6}$ and $\sim 10^{-9}$. Anomalous τ magnetic moments arising within seesaw models were explored in Ref. [60], where values of orders 10^{-8} and 10^{-9} were respectively determined for the type I and type III versions of seesaw. It was found in Ref. [61] that different scenarios allow a spin-0 unparticle to produce contributions to a_τ in the wide range of values $\sim 10^{-10} - \sim 10^{-6}$, though the authors of this reference noticed and emphasized that certain scenarios could generate even larger contributions than that of the Standard Model. Ref. [62] includes a calculation of the contributions from scalar leptoquark interactions to

the τ anomalous magnetic moment, with values between $\sim 10^{-9}$ and $\sim 10^{-8}$.

IV. ELECTRIC DIPOLE MOMENTS AT ONE LOOP

In this section we turn our attention to the diagonal electric dipole moments, whose general structure is the one given in Eq. (4). Since the assumption that $v_{j\alpha} \neq a_{j\alpha}$ is a necessary requirement to have nonzero electric dipole moments, we find convenient to introduce the difference $\Delta_{j\alpha} \equiv a_{j\alpha} - v_{j\alpha}$. Furthermore, because CP violation is also necessary, we assume that $\Delta_{j\alpha}$ and $v_{j\alpha}$ are complex quantities. Analogously to what we did with the anomalous magnetic moments, we denote each heavy-neutrino contribution to the electric dipole moment as $d_{j\alpha}$, meaning that $d_\alpha = \sum_j d_{j\alpha}$, with any of such contributions expressed as

$$d_{j\alpha} \approx \frac{1}{(16\pi)^2} \sin \phi_{j\alpha} \frac{|\Delta_{j\alpha}|}{m_j} \left[\frac{5}{2} - \frac{1}{2} \frac{m_\alpha^2}{m_j^2} + \epsilon_j \left(\frac{13}{3} - \frac{5}{3} \frac{m_\alpha^2}{m_j^2} \right) \right] \cdot e. \quad (12)$$

To write down this equation, we have taken $2|v_{j\alpha}| \sim 1$. The angle $\phi_{j\alpha}$ is a phase difference of the complex phases of $v_{j\alpha}$ and $\Delta_{j\alpha}$. Note that this phase difference can be conveniently chosen to appropriately adjust the sign of the contribution. Of course, if CP were preserved then $\phi_{j\alpha} = 0$, which would eliminate the $d_{j\alpha}$ contribution. The contributions $d_{j\alpha}$ decouple for $m_j \rightarrow \infty$, which is quite explicit in Eq. (12).

So far, experiments have not observed electric dipole moments of elementary particles, and this lack of measurements has been translated into upper bounds. The electric dipole moment of the electron has received special attention among all electric dipoles of elementary particles, being the one which has been most stringently bounded. Experiments with thallium atoms and ytterbium fluoride molecules achieved upper bounds of order $10^{-27} e \cdot \text{cm}$ on $|d_e|$ [63–65]. The current champion, however, is the upper bound recently established by the ACME Collaboration [39], which reached an important improvement finding $|d_e| < 8.7 \times 10^{-29} e \cdot \text{cm}$, at 90% C.L. The huge difference, of 15 orders of magnitude, between the Standard Model contribution [41] and the current experimental sensitivity has motivated the introduction of new CP -violating physics, pursuing less suppressed values for this observable. Particularly, the authors of Ref. [66] asserted that in presence of Majorana neutrinos, a two-loop contribution to the electric dipole moment of the electron is produced. Other investigations concerning electric dipole moments and massive neutrinos were carried out in Refs. [67, 68]. As we did with the anomalous magnetic moments, we take $d_\alpha = \sum_j d_{j\alpha} \approx 3 d_{N\alpha}$. Our discussion of the last section led us to conclude

that ϵ_N lies between 10^{-13} and 10^{-10} . On the other hand, note that for heavy masses m_N within the range 0.5 TeV–7 TeV, we have ratios $m_e^2/m_N^2 \sim 10^{-15} - 10^{-12}$. Clearly, both contributions are well below the 5/2 term, which allows us to establish

$$|\Delta_{Ne}| \lesssim 2.53 \times 10^{-13} \left(\frac{m_N}{1 \text{ GeV}} \right). \quad (13)$$

From the last equation we note that a heavy neutrino with mass $m_N \sim 10^2$ GeV shall yield a difference $|\Delta_{je}|$ of order 10^{-11} , while a difference of order 10^{-10} can be achieved with a mass $m_N \sim 10^3$ GeV. Recall that in the discussion of the anomalous magnetic moments we used that $|\Delta_{j\alpha}| \ll m_\alpha/m_j$. As we just saw, while $|\Delta_{j\alpha}| \sim 10^{-11} - 10^{-10}$, we have $m_e/m_N \sim 10^{-8} - 10^{-7}$, which supports our previous assumption.

V. CONCLUSIONS

In the light of the confirmation that the phenomenon of neutrino oscillations exist, including neutrino mixing, neutrino mass and CP violation, in this paper we have explored a model in which heavy neutrinos and a heavy W' gauge boson, both originated in some high energy formulation, couple with charged leptons from the Standard Model in a set of general charged currents characterized by vector and axial terms that differ of each other, though in a small amount. We have shown that, as long as such difference is present, CP is violated by these charged currents, and neutrinos are massive, nonzero contributions from such charged currents to electric dipole moments of charged leptons arise at one loop. We have also performed an analysis of the contributions to anomalous magnetic moments, which, in this model, decouple and require that the masses of the heavy neutrinos constitute

a quasi-degenerate spectrum. Using the current difference among experiment and the Standard Model prediction of the anomalous magnetic moments of the electron and the muon, we have determined a narrow allowed region for a parameter ϵ_N , characterizing the difference among the masses of the neutrinos and the W' boson. We have used this information to provide an estimation of the anomalous magnetic moment of the tau lepton, which turned out to be within the range $\sim 10^{-10} - 10^{-8}$. Taking a small parameter ϵ_N , we have used the stringent upper bound on the electron electric dipole moment, recently reported by the ACME Collaboration, to estimate that the difference between the vector and axial parts of the charged currents characterizing the electron must be $\lesssim 10^{-11} - 10^{-10}$, for heavy neutrino masses in the range 0.5 TeV – 7 TeV.

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Appendix A: Exact analytic expressions of factors in the magnetic and electric moments

In Eqs. (3) to (6) we gave the general form of the contributions to magnetic and electric moments of charged leptons. These expressions are written in terms of mass-dependent factors $M_{j\alpha}^A$, $M_{j\alpha}^V$, $D_{j\alpha}$, $M_{j,\alpha\beta}^A$, $M_{j,\alpha\beta}^V$, $D_{j,\alpha\beta}$, and $\overline{D}_{j,\alpha\beta}$, whose exact analytic expressions are provided in this Appendix. Using the standard definitions

$$B_0(p^2, m_0^2, m_1^2) = \frac{(2\pi\mu)^{4-D}}{i\pi^2} \int d^D k \frac{1}{[k^2 - m_0^2][(k+p)^2 - m_1^2]}, \quad (A1)$$

$$C_0(q^2, p^2, (q-p)^2, m_0^2, m_1^2, m_2^2) = \frac{(2\pi\mu)^{4-D}}{i\pi^2} \int d^D k \frac{1}{[k^2 - m_0^2][(k+q)^2 - m_1^2][(k+q-p)^2 - m_2^2]}, \quad (A2)$$

we have the following expressions:

$$\begin{aligned}
M_{j\alpha}^A = & \frac{1}{2(16\pi)^2 m_\alpha^2 m_{W'}^2} \left[2m_\alpha^2 ((m_\alpha + m_j)^2 + 2m_{W'}^2) + 2(m_{W'}^2 - m_j^2) ((m_\alpha + m_j)^2 + 2m_{W'}^2) B_0(0, m_j^2, m_{W'}^2) \right. \\
& + (m_\alpha^4 + m_\alpha^2 (13m_{W'}^2 - 2m_j^2) + 6m_\alpha m_j (m_j^2 + m_{W'}^2) + 5(m_j^4 + m_j^2 m_{W'}^2 - 2m_{W'}^4) \\
& - 2m_\alpha^3 m_j) B_0(m_\alpha^2, m_j^2, m_{W'}^2) + (-m_\alpha^4 + m_\alpha^2 (4m_j^2 - 15m_{W'}^2) - 2m_\alpha m_j (m_j^2 + 5m_{W'}^2) \\
& - 3(m_j^4 + m_j^2 m_{W'}^2 - 2m_{W'}^4) + 2m_\alpha^3 m_j) B_0(0, m_{W'}^2, m_{W'}^2) + ((m_\alpha + m_j)^2 - m_{W'}^2) (m_\alpha^4 - 3m_j^4 + 4m_\alpha m_j^3 \\
& \left. + m_j^2 (2m_\alpha^2 - 3m_{W'}^2) + 2m_\alpha m_j (m_{W'}^2 - 2m_\alpha^2) + 6m_{W'}^4 - 11m_\alpha^2 m_{W'}^2) C_0(m_\alpha^2, m_\alpha^2, 0, m_{W'}^2, m_j^2, m_{W'}^2) \right], \quad (A3)
\end{aligned}$$

$$\begin{aligned}
M_{j\alpha}^V = & \frac{1}{2(16\pi^2) m_\alpha^2 m_{W'}^2} \left[2m_\alpha^2 ((m_\alpha - m_j)^2 + 2m_{W'}^2) + 2(m_{W'}^2 - m_j^2) ((m_\alpha - m_j)^2 + 2m_{W'}^2) B_0(0, m_j^2, m_{W'}^2) \right. \\
& + (m_\alpha^4 + m_\alpha^2 (13m_{W'}^2 - 2m_j^2) - 6m_\alpha m_j (m_j^2 + m_{W'}^2) + 5(m_j^4 + m_j^2 m_{W'}^2 - 2m_{W'}^4) \\
& + 2m_\alpha^3 m_j) B_0(m_\alpha^2, m_j^2, m_{W'}^2) + (-m_\alpha^4 + m_\alpha^2 (4m_j^2 - 15m_{W'}^2) + 2m_\alpha m_j (m_j^2 + 5m_{W'}^2) \\
& - 3(m_j^4 + m_j^2 m_{W'}^2 - 2m_{W'}^4) - 2m_\alpha^3 m_j) B_0(0, m_{W'}^2, m_{W'}^2) + ((m_\alpha - m_j)^2 - m_{W'}^2) (m_\alpha^4 - 3m_j^4 - 4m_\alpha m_j^3 \\
& \left. + m_j^2 (2m_\alpha^2 - 3m_{W'}^2) - 2m_\alpha m_j (m_{W'}^2 - 2m_\alpha^2) + 6m_{W'}^4 - 11m_\alpha^2 m_{W'}^2) C_0(m_\alpha^2, m_\alpha^2, 0, m_{W'}^2, m_j^2, m_{W'}^2) \right], \quad (A4)
\end{aligned}$$

$$\begin{aligned}
D_{j\alpha} = & \frac{1}{(16\pi)^2} \frac{m_j}{m_\alpha^2 m_{W'}^2} \left[(m_\alpha^2 - m_j^2 + 4m_{W'}^2) (B_0(m_\alpha^2, m_j^2, m_{W'}^2) - B_0(0, m_{W'}^2, m_{W'}^2)) \right. \\
& \left. + (m_\alpha^4 + m_j^4 - m_j^2 (2m_\alpha^2 + 5m_{W'}^2) + 4m_{W'}^4 - 3m_\alpha^2 m_{W'}^2) C_0(m_\alpha^2, m_\alpha^2, 0, m_{W'}^2, m_j^2, m_{W'}^2) \right], \quad (A5)
\end{aligned}$$

$$\begin{aligned}
M_{j,\alpha\beta}^A = & \frac{1}{(16\pi)^2} \frac{1}{m_\alpha m_\beta m_{W'}^2 (m_\alpha^2 - m_\beta^2)^2} \left[(m_\alpha^2 - m_\beta^2)^2 (m_j^2 - m_{W'}^2) ((m_\alpha + m_j)(m_\beta + m_j) + 2m_{W'}^2) B_0(0, m_j^2, m_{W'}^2) \right. \\
& - m_\alpha (m_\alpha^2 - m_\beta^2) (m_j^4 (m_\alpha + 2m_\beta) + m_j^3 m_\alpha (m_\alpha + m_\beta) + m_j^2 (m_{W'}^2 (m_\alpha + 2m_\beta) \\
& - m_\beta (m_\alpha m_\beta + m_\alpha^2 + 2m_\beta^2)) - m_j (m_\alpha + m_\beta) (m_\alpha m_\beta^2 + m_{W'}^2 (m_\alpha - 6m_\beta)) + m_\alpha^2 m_\beta^3 \\
& - 2m_{W'}^4 (m_\alpha + 2m_\beta) + m_\beta m_{W'}^2 (m_\alpha + 2m_\beta) (3m_\alpha + 2m_\beta)) B_0(m_\alpha^2, m_j^2, m_{W'}^2) \\
& + m_\beta (m_\alpha^2 - m_\beta^2) (m_j^4 (2m_\alpha + m_\beta) + m_j^3 m_\beta (m_\alpha + m_\beta) + m_j^2 (m_{W'}^2 (2m_\alpha + m_\beta) \\
& - m_\alpha (m_\alpha m_\beta + 2m_\alpha^2 + m_\beta^2)) - m_j (m_\alpha + m_\beta) (m_\alpha^2 m_\beta + m_{W'}^2 (m_\beta - 6m_\alpha)) + m_\alpha^3 m_\beta^2 \\
& - 2m_{W'}^4 (2m_\alpha + m_\beta) + m_\alpha m_{W'}^2 (2m_\alpha + m_\beta) (2m_\alpha + 3m_\beta)) B_0(m_\beta^2, m_j^2, m_{W'}^2) \\
& - 2m_\alpha m_\beta m_{W'}^2 (m_\alpha^2 - m_\beta^2)^2 (-3m_j (m_\alpha + m_\beta) + m_j^2 - 3m_\alpha m_\beta - 2m_\alpha^2 - 2m_\beta^2 \\
& \left. + 2m_{W'}^2) C_0(m_\alpha^2, m_\beta^2, 0, m_{W'}^2, m_j^2, m_{W'}^2) - m_\alpha m_\beta (m_\alpha^2 - m_\beta^2)^2 ((m_j + m_\alpha)(m_j + m_\beta) + 2m_{W'}^2) \right], \quad (A6)
\end{aligned}$$

$$\begin{aligned}
M_{j,\alpha\beta}^V = & \frac{1}{(16\pi)^2} \frac{1}{m_\alpha m_\beta m_{W'}^2 (m_\alpha^2 - m_\beta^2)^2} \left[(m_\alpha^2 - m_\beta^2)^2 (m_j^2 - m_{W'}^2) ((m_j - m_\alpha)(m_j - m_\beta) + 2m_{W'}^2) B_0(0, m_j^2, m_{W'}^2) \right. \\
& - m_\alpha (m_\alpha^2 - m_\beta^2) (m_j^4 (m_\alpha + 2m_\beta) - m_j^3 m_\alpha (m_\alpha + m_\beta) + m_j^2 (m_{W'}^2 (m_\alpha + 2m_\beta) \\
& - m_\beta (m_\alpha m_\beta + m_\alpha^2 + 2m_\beta^2)) + m_j (m_\alpha + m_\beta) (m_\alpha m_\beta^2 + m_{W'}^2 (m_\alpha - 6m_\beta)) + m_\alpha^2 m_\beta^3 \\
& - 2m_{W'}^4 (m_\alpha + 2m_\beta) + m_\beta m_{W'}^2 (m_\alpha + 2m_\beta) (3m_\alpha + 2m_\beta)) B_0(m_\alpha^2, m_j^2, m_{W'}^2) \\
& + m_\beta (m_\alpha^2 - m_\beta^2) (m_j^4 (2m_\alpha + m_\beta) - m_j^3 m_\beta (m_\alpha + m_\beta) + m_j^2 (m_{W'}^2 (2m_\alpha + m_\beta) \\
& - m_\alpha (m_\alpha m_\beta + 2m_\alpha^2 + m_\beta^2)) + m_j (m_\alpha + m_\beta) (m_\alpha^2 m_\beta + m_{W'}^2 (m_\beta - 6m_\alpha)) + m_\alpha^3 m_\beta^2 \\
& - 2m_{W'}^4 (2m_\alpha + m_\beta) + m_\alpha m_{W'}^2 (2m_\alpha + m_\beta) (2m_\alpha + 3m_\beta)) B_0(m_\beta^2, m_j^2, m_{W'}^2) \\
& - 2m_\alpha m_\beta m_{W'}^2 (m_\alpha^2 - m_\beta^2)^2 (3m_j (m_\alpha + m_\beta) + m_j^2 - 3m_\alpha m_\beta - 2m_\alpha^2 - 2m_\beta^2 \\
& \left. + 2m_{W'}^2) C_0(m_\alpha^2, m_\beta^2, 0, m_{W'}^2, m_j^2, m_{W'}^2) - m_\alpha m_\beta (m_\alpha^2 - m_\beta^2)^2 ((m_j - m_\alpha)(m_j - m_\beta) + 2m_{W'}^2) \right], \quad (A7)
\end{aligned}$$

$$\begin{aligned}
D_{j,\alpha\beta} = & \frac{1}{(16\pi)^2} \frac{m_\alpha + m_\beta}{m_\alpha m_\beta m_{W'}^2 (m_\alpha^2 - m_\beta^2)^2} \left[(m_j^2 - m_{W'}^2) (m_\beta^2 - m_\alpha^2) ((m_j - m_\alpha)(m_j + m_\beta) + 2m_{W'}^2) B_0(0, m_j^2, m_{W'}^2) \right. \\
& + m_\alpha (m_j^4 (m_\alpha - 2m_\beta) + m_j^3 m_\alpha (m_\beta - m_\alpha) + m_j^2 (m_\beta (-m_\alpha m_\beta + m_\alpha^2 + 2m_\beta^2) + m_{W'}^2 (m_\alpha - 2m_\beta)) \\
& + m_j (m_\alpha - m_\beta) (m_\alpha m_\beta^2 + m_{W'}^2 (m_\alpha + 6m_\beta)) - m_\alpha^2 m_\beta^3 - 2m_{W'}^4 (m_\alpha - 2m_\beta) \\
& - m_\beta m_{W'}^2 (m_\alpha - 2m_\beta) (3m_\alpha - 2m_\beta)) B_0(m_\alpha^2, m_j^2, m_{W'}^2) \\
& + m_\beta (m_j^4 (2m_\alpha - m_\beta) + m_j^3 m_\beta (m_\alpha - m_\beta) + m_j^2 (m_{W'}^2 (2m_\alpha - m_\beta) - m_\alpha (-m_\alpha m_\beta + 2m_\alpha^2 + m_\beta^2)) \\
& - m_j (m_\alpha - m_\beta) (m_\alpha^2 m_\beta + m_{W'}^2 (6m_\alpha + m_\beta)) + m_\alpha^3 m_\beta^2 + 2m_{W'}^4 (m_\beta - 2m_\alpha) \\
& + m_\alpha m_{W'}^2 (-8m_\alpha m_\beta + 4m_\alpha^2 + 3m_\beta^2)) B_0(m_\beta^2, m_j^2, m_{W'}^2) \\
& - 2m_\alpha m_\beta m_{W'}^2 (m_\alpha^2 - m_\beta^2) (3m_j (m_\alpha - m_\beta) + m_j^2 + 3m_\alpha m_\beta - 2m_\alpha^2 - 2m_\beta^2) \\
& \left. + 2m_{W'}^2) C_0(m_\alpha^2, m_\beta^2, 0, m_{W'}^2, m_j^2, m_{W'}^2) - m_\alpha m_\beta (m_\alpha^2 - m_\beta^2) ((m_j - m_\alpha)(m_j + m_\beta) + 2m_{W'}^2) \right], \quad (A8)
\end{aligned}$$

$$\begin{aligned}
\overline{D}_{j,\alpha\beta} = & \frac{1}{(16\pi)^2} \frac{m_\alpha + m_\beta}{m_\alpha m_\beta m_{W'}^2 (m_\alpha^2 - m_\beta^2)^2} \left[(m_j^2 - m_{W'}^2) (m_\beta^2 - m_\alpha^2) ((m_j + m_\alpha)(m_j - m_\beta) + 2m_{W'}^2) B_0(0, m_j^2, m_{W'}^2) \right. \\
& + m_\alpha (m_j^4 (m_\alpha - 2m_\beta) + m_j^3 m_\alpha (m_\alpha - m_\beta) + m_j^2 (m_\beta (-m_\alpha m_\beta + m_\alpha^2 + 2m_\beta^2) + m_{W'}^2 (m_\alpha - 2m_\beta)) \\
& - m_j (m_\alpha - m_\beta) (m_\alpha m_\beta^2 + m_{W'}^2 (m_\alpha + 6m_\beta)) - m_\alpha^2 m_\beta^3 - 2m_{W'}^4 (m_\alpha - 2m_\beta) \\
& - m_\beta m_{W'}^2 (m_\alpha - 2m_\beta) (3m_\alpha - 2m_\beta)) B_0(m_\alpha^2, m_j^2, m_{W'}^2) \\
& + m_\beta (m_j^4 (2m_\alpha - m_\beta) + m_j^3 m_\beta (m_\beta - m_\alpha) + m_j^2 (m_{W'}^2 (2m_\alpha - m_\beta) - m_\alpha (-m_\alpha m_\beta + 2m_\alpha^2 + m_\beta^2)) \\
& + m_j (m_\alpha - m_\beta) (m_\alpha^2 m_\beta + m_{W'}^2 (6m_\alpha + m_\beta)) + m_\alpha^3 m_\beta^2 + 2m_{W'}^4 (m_\beta - 2m_\alpha) \\
& + m_\alpha m_{W'}^2 (-8m_\alpha m_\beta + 4m_\alpha^2 + 3m_\beta^2)) B_0(m_\beta^2, m_j^2, m_{W'}^2) \\
& - 2m_\alpha m_\beta m_{W'}^2 (m_\alpha^2 - m_\beta^2) (3m_\beta (m_j + m_\alpha) - 3m_j m_\alpha + m_j^2 - 2m_\alpha^2 - 2m_\beta^2) \\
& \left. + 2m_{W'}^2) C_0(m_\alpha^2, m_\beta^2, 0, m_{W'}^2, m_j^2, m_{W'}^2) - m_\alpha m_\beta (m_\alpha^2 - m_\beta^2) ((m_j + m_\alpha)(m_j - m_\beta) + 2m_{W'}^2) \right]. \quad (A9)
\end{aligned}$$

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